

The Higgs boson mass and SUSY spectra in 10D SYM theory with magnetized extra dimensions

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Abstract

We study the Higgs boson mass and the spectrum of supersymmetric (SUSY) particles in the well-motivated particle physics model derived from a ten-dimensional supersymmetric Yang-Mills theory compactified on three factorizable tori with magnetic fluxes. This model was proposed in a previous work, where the flavor structures of the standard model including the realistic Yukawa hierarchies are obtained from non-hierarchical input parameters on the magnetized background. Assuming moduli- and anomaly-mediated contributions dominate the soft SUSY breaking terms, we study the precise SUSY spectra and analyze the Higgs boson mass in this mode, which are compared with the latest experimental data.

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1 Introduction

The discovery of a Higgs boson at the Large Hadron Collider (LHC) [1, 2] provides the last complement of the standard model (SM), and it is getting recognized as the successful and reliable theory more than before. However, various phenomenological (or other) facts clearly indicate the presence of new physics beyond the SM. Supersymmetry (SUSY) and extra dimensional space are the well-known great candidates for that and diverse models has been constructed with them. Now, the discovery of the Higgs boson enables us to further study or verify such models including the SUSY spectra and dynamics of the extra dimensional space.

The complicated structures, e.g., product gauge groups, three generations and hierarchical Yukawa couplings, are some of the most well-known mysteries of the SM. In this paper with the motive, we consider higher-dimensional supersymmetric Yang-Mills (SYM) theories compactified on simple factorizable tori with magnetic fluxes. The magnetic fluxes on the tori induce the four-dimensional (4D) chiral spectra [3, 4], and furthermore, that is capable of originating the flavors of the SM. The degenerate zero-modes are induced by the magnetic fluxes with a certain degeneracy corresponding to the flux magnitude which is quantized due to the Dirac's argument, and their wavefunctions are localized at different points of the magnetized tori. As a result, Yukawa coupling constants in the 4D effective theory, which are given by overlap integrals of such localized wavefunctions, can be hierarchical among the degenerate zero-modes, that are the generations.

Toroidal compactifications with magnetic fluxes are quite attractive to derive the SM from the higher-dimensional SYM theories. The SUSY theories in higher-dimensional spacetime have $\mathcal{N} = 2, 3$ or 4 SUSY counted by the 4D supercharges depending on the structure of background as well as the dimensionality of spacetime. The magnetic fluxes generally break the SUSY and careful analyses are required to construct MSSM-like models preserving the $\mathcal{N} = 1$ SUSY, because the configuration of fluxes determine not only the number of SUSY preserved but also the almost everything mentioned above. In the 4D $\mathcal{N} = 1$ superspace, from such a perspective, we proposed a systematic way of dimensional reduction of the ten-dimensional (10D) SYM theories compactified on three factorizable tori with magnetic fluxes [5]. Thanks to that, a MSSM-like model preserving the $\mathcal{N} = 1$ SUSY was constructed [6], where the magnetic fluxes originated the flavor structures of the SM particles and even of their superpartners by assuming some typical mediation mechanisms of SUSY breaking. Then the observational consistencies with respect to the SUSY flavor violations were also studied in Ref [6].

In this paper, we study the low-energy phenomenology of the model proposed by Ref [6], especially focusing on the Higgs boson mass and the SUSY spectrum precisely with the latest experimental data, and those two issues greatly correlate to each other. In the model, the background fluxes are already fixed to realize the SM flavor structures, then there are only a few remaining parameters which appear mainly from the degrees of freedom governing the (SUSY) dynamics of extra dimensional space, i.e., the moduli superfields, determining the SUSY spectrum in the MSSM sector. We will study the parameter dependence of the Higgs boson mass via the SUSY spectrum.

The organization of this paper is as follows. In Section 2, we give a brief review of the magnetized model proposed in Ref. [6] including the way of dimensional reduction with the superfield

description [5]. In Section 3, we estimate the masses and mixing angles of the quarks and the leptons. We can calculate the Yukawa coupling constants depending on some parameters in the model. Such estimations have already been done, but we improve them especially taking the Higgs boson mass into account. On the background, we study the Higgs boson mass and the SUSY spectrum in Section 4. We assume that the moduli- and anomaly-mediated contributions dominate the soft SUSY breaking terms in the MSSM sector, which are typical mediators in higher-dimensional spacetime, and parameterize them by the auxiliary F -components of the corresponding moduli and the compensator supermultiplets, respectively. Accordingly, we calculate the soft SUSY breaking parameters, the SUSY spectrum, and furthermore, the Higgs boson mass. Finally, Section 5 is devoted to conclusions and discussions.

2 Review of the magnetized model

We give a review of the phenomenological model proposed in Ref. [6]. The model was derived starting from the 10D $U(8)$ SYM theory compactified on a (factorizable) product of 4D Minkowski space and three two-dimensional (2D) tori, $R^{1,3} \times \prod_{i=1}^3 (T^2)_i$. The action is given by

$$S = \int d^{10}x \sqrt{-G} \left\{ -\frac{1}{4g^2} \text{tr} (F^{MN} F_{MN}) + \frac{i}{2g^2} \text{tr} (\bar{\lambda} \Gamma^M D_M \lambda) \right\}, \quad (1)$$

that contains the 10D vector field A_M ($M = 0, 1, \dots, 9$) and the 10D Majorana-Weyl spinor field λ . For each $i = 1, 2, 3$, the metric of 2D torus $(T^2)_i$ is expressed as

$$g^{(i)} = (2\pi R^{(i)})^2 \begin{pmatrix} 1 & \text{Re}\tau^{(i)} \\ \text{Re}\tau^{(i)} & |\tau^{(i)}|^2 \end{pmatrix},$$

where the real parameters $R^{(i)}$ correspond to typical sizes of the three torus and the complex parameters $\tau^{(i)}$ to the complex structures ($i = 1, 2, 3$), and those are all contained in the 10D metric G_{MN} whose determinant is described by G in the action.

We will introduce magnetic fluxes on each of the factorizable tori and derive the 4D effective action. We have an extremely useful superfield description for such dimensional reductions. (The details are given in Ref. [5].) First, we define the complex coordinates z_i of the i -th 2D torus $(T^2)_i$ as

$$z^i \equiv \frac{1}{2} (x^{2+2i} + \tau^{(i)} x^{3+2i}),$$

and its metric $h_{i\bar{j}} = 2(2\pi R^{(i)})^2 \delta_{i\bar{j}} = \delta_{i\bar{j}} e_i^{\mathbf{i}} e_{\bar{j}}^{\mathbf{j}}$ with which the line element in 6D extra space is given by $ds_{6D}^2 = 2h_{i\bar{j}} dz^i d\bar{z}^{\bar{j}}$ where $\bar{z}^{\bar{i}}$ is the complex conjugate to z^i . The 10D vector field A_M is decomposed into the 4D vector field A_μ and the others, and we define three complex fields,

$$A_i \equiv -\frac{1}{\text{Im}\tau^{(i)}} (\bar{\tau}^{(i)} A_{2+2i} - A_{3+2i}).$$

The 10D Majorana-Weyl spinor field λ is also decomposed into four 4D Weyl spinors, $\lambda_0 = \lambda_{+++}$, $\lambda_1 = \lambda_{+--}$, $\lambda_2 = \lambda_{-+-}$ and $\lambda_3 = \lambda_{--+}$. The subscripts \pm represent the chirality on

the three tori and other combinations of them, λ_{---} , λ_{-++} , λ_{+-+} and λ_{++-} do not appear because of the 10D Majorana-Weyl condition.

These form a vector multiplet $\{A_\mu, \lambda_0\}$ and three chiral multiplets $\{A_i, \lambda_i\}$ under a 4D $\mathcal{N} = 1$ SUSY, which are assigned to a 4D $\mathcal{N} = 1$ vector superfield V and three chiral superfields ϕ_i respectively as follows,

$$\begin{aligned} V &\equiv -\theta\sigma^\mu\bar{\theta}A_\mu + i\bar{\theta}\bar{\theta}\theta\lambda_0 - i\theta\theta\bar{\theta}\bar{\lambda}_0 + \frac{1}{2}\theta\theta\bar{\theta}\bar{\theta}D, \\ \phi_i &\equiv \frac{1}{\sqrt{2}}A_i + \sqrt{2}\theta\lambda_i + \theta\theta F_i, \end{aligned} \quad (2)$$

where the θ denotes the Grassmann coordinate of the $\mathcal{N} = 1$ superspace. These superfields lead us to rewrite the SYM action (1) in the 4D $\mathcal{N} = 1$ superspace as [5, 7]

$$S = \int d^{10}x\sqrt{-G} \left[\int d^4\theta \mathcal{K} + \left\{ \int d^2\theta \left(\frac{1}{4g^2} \mathcal{W}^\alpha \mathcal{W}_\alpha \right) + \text{h.c.} \right\} \right], \quad (3)$$

with the following functions of superfields,

$$\begin{aligned} \mathcal{K} &= \frac{2}{g^2} h^{i\bar{j}} \text{Tr} \left[\left(\sqrt{2}\bar{\partial}_{\bar{i}} + \bar{\phi}_{\bar{i}} \right) e^{-V} \left(-\sqrt{2}\partial_j + \phi_j \right) e^V + \bar{\partial}_{\bar{i}} e^{-V} \partial_j e^V \right] + \mathcal{K}_{\text{WZW}}, \\ \mathcal{W} &= \frac{1}{g^2} \epsilon^{ijk} e_i^i e_j^j e_k^k \text{tr} \left[\sqrt{2}\phi_i \left(\partial_j \phi_k - \frac{1}{3\sqrt{2}} [\phi_j, \phi_k] \right) \right], \\ \mathcal{W}_\alpha &= -\frac{1}{4} \bar{D}\bar{D}e^{-V} D_\alpha e^V, \end{aligned} \quad (4)$$

where ∂_i is a derivative with respect to the complex coordinate z_i , and h_{ij} and e_i^i are the metric and vielbein of the three tori. The terms \mathcal{K}_{WZW} vanish in the Wess-Zumino gauge. In the definition of \mathcal{W}_α , D_α and \bar{D}_α denote the (super)covariant derivatives. The 10D SYM action with the superfield description has the $\mathcal{N} = 4$ SUSY, and a $\mathcal{N} = 1$ SUSY, which is a part of that, becomes manifest. The superfields V and ϕ_i contain the auxiliary fields D and F_i respectively, and it is obvious that the $\mathcal{N} = 1$ SUSY is preserved if their vacuum expectation values (VEVs) vanish.

We consider the following magnetized background,

$$\langle A_i \rangle = \frac{\pi}{\text{Im}\tau^{(i)}} \left(M^{(i)} \bar{z}_i + \bar{\zeta}_i \right), \quad (5)$$

where $M^{(i)}$ and ζ_i are matrices of the number of magnetic fluxes and of the Wilson-lines on $(T^2)_i$. The SUSY preserving condition $\langle D \rangle = 0$ requires (The $(1,1)$ form fluxes (5) automatically satisfy $\langle F_i \rangle = 0$.)

$$\frac{1}{\mathcal{A}^{(1)}} M^{(1)} + \frac{1}{\mathcal{A}^{(2)}} M^{(2)} + \frac{1}{\mathcal{A}^{(3)}} M^{(3)} = 0, \quad (6)$$

where $\mathcal{A}^{(i)} = (2\pi R^{(i)})^2 \text{Im}\tau^{(i)}$ expresses the area of $(T^2)_i$. In our model, the magnetic fluxes and Wilson-line matrices are encoded in the following (8×8) diagonal matrices whose rows

and columns cover the space of the $U(8)$ gauge group,

$$M^{(i)} = \begin{pmatrix} M_C^{(i)} \mathbf{1}_4 & 0 & 0 \\ 0 & M_L^{(i)} \mathbf{1}_2 & 0 \\ 0 & 0 & M_R^{(i)} \mathbf{1}_2 \end{pmatrix}, \quad \zeta^{(i)} = \begin{pmatrix} \zeta_C^{(i)} \mathbf{1}_3 & 0 & 0 & 0 & 0 \\ 0 & \zeta_{C'}^{(i)} & 0 & 0 & 0 \\ 0 & 0 & \zeta_L^{(i)} \mathbf{1}_2 & 0 & 0 \\ 0 & 0 & 0 & \zeta_{R'}^{(i)} & 0 \\ 0 & 0 & 0 & 0 & \zeta_{R''}^{(i)} \end{pmatrix}. \quad (7)$$

The explicit values of $M_a^{(i)}$ and $\zeta_a^{(i)}$ are determined phenomenologically which are shown later for each $a = C, C', L, R', R''$. The flux matrices $M^{(i)}$ must have the three-block diagonal forms because that is strongly constrained by the SUSY condition and the diagonal components $M_a^{(i)}$ are integers due to the Dirac's quantization condition. They break the $U(8)$ gauge group down to $U(4)_C \times U(2)_L \times U(2)_R$, the Pati-Salam gauge group, and that is further broken down by the Wilson-lines close to the SM gauge group. There are five unbroken $U(1)$ s and a linear combination of them will be the $U(1)$ hypercharge. The gauge bosons of the other four $U(1)$ s than that of the $U(1)_Y$ are assumed to become massive due to some UV physics (e.g., the Green-Schwartz mechanism [8]) and decouple from the physics below the compactification scale. We use the indices, $a, b = C, C', L, R', R''$ for these remaining subgroups and ϕ_i^{ab} represents a bifundamental representation (N_a, \bar{N}_b) of $U(N_a) \times U(N_b)$, included in ϕ_i which is the adjoint representation of $U(8)$ for each $i = 1, 2, 3$. These notations are similarly adopted for V which is also the adjoint representation of $U(8)$.

From now on, we are focusing on only the zero-modes of V^{ab} and ϕ_i^{ab} to construct the phenomenological model which is being implicit, and then we use the same notation for the zero-mode as the one for the corresponding 10D fields. The degenerate zero-modes appear depending on the flux configuration and their wavefunctions can be obtained solving the zero-mode equations [4, 5]. In the following, we summarize the results briefly.

First, the diagonal parts of the vector superfield V^{aa} correspond to the SM gauge fields. They feel no magnetic flux and have a flat wavefunction on the torus. Those of the chiral superfields ϕ_i^{aa} also remain as massless vector-like exotics or notorious open string moduli. In the phenomenological model building, some prescriptions are required to treat these zero-modes ϕ_i^{aa} if exist as we will mention it later.

Next, the off-diagonals $V^{ab} (a \neq b)$ have heavy masses in response to the partial breaking of the $U(8)$ gauge group, and they have no effect on the phenomenologies at the low-energy below the compactification scale. As for the bifundamentals ϕ_i^{ab} , they carry the matter fields and the Higgs fields of the MSSM, which are the most important phenomenological ingredients. The bifundamentals ϕ_i^{ab} feel the magnetic fluxes $M_{ab}^{(j)} \equiv M_a^{(j)} - M_b^{(j)}$ on the torus $(T^2)_j$. In the case with $i = j$, degenerate zero-modes arise if and only if $M_{ab}^{(j)} > 0$, whose degeneracy is also given by $M_{ab}^{(j)}$, while the conjugate representations ϕ_i^{ba} are eliminated because they feel the negative magnetic fluxes $M_{ba}^{(j)} < 0$. On the other hand, in the case with $i \neq j$, negative magnetic fluxes $M_{ab}^{(j)} < 0$ yield the degenerate zero-modes with the degeneracy $|M_{ab}^{(j)}|$, and their conjugates are projected out because of the positive magnetic fluxes $M_{ab}^{(j)} > 0$. Consequently, the magnetic fluxes cause a kind of projection generating a 4D chiral spectra

with the degeneracies identified as generations. If there are vanishing magnetic fluxes $M_{ab}^{(j)} = M_{ba}^{(j)} = 0$, the projection does not occur and both of the single zero-modes ϕ_i^{ab} and ϕ_i^{ba} remain simultaneously with flat wavefunctions regardless of whether $i = j$ or not. We can identify the degenerate zero-modes as the generations of the MSSM matters.

Furthermore, the magnetic fluxes give the Gaussian-like profiles to the zero-mode wavefunctions, and each degenerate zero-mode is localized at the different point on the magnetized torus from each other. The Yukawa coupling constants in the 4D effective action are given by overlap integrals of these Gaussian-like wavefunctions. The magnitude of the overlaps determines the values of the Yukawa coupling constants, while the Wilson-lines can shift the peak position of each Gaussian differently in general, that is, they can control the value of the Yukawa coupling constants hierarchically.

Now, we show the phenomenologically specific flux configuration,

$$\begin{aligned} (M_C^{(1)}, M_L^{(1)}, M_R^{(1)}) &= (0, 3, -3), \\ (M_C^{(2)}, M_L^{(2)}, M_R^{(2)}) &= (0, -1, 0), \\ (M_C^{(3)}, M_L^{(3)}, M_R^{(3)}) &= (0, 0, 1), \end{aligned} \tag{8}$$

on which magnetized background, the SUSY preserving condition (6) is satisfied if

$$\mathcal{A}^{(1)}/\mathcal{A}^{(2)} = \mathcal{A}^{(1)}/\mathcal{A}^{(3)} = 3. \tag{9}$$

This flux configuration is the unique one to produce the three-generation MSSM-like model preserving the $\mathcal{N} = 1$ SUSY unless we consider more complicated magnetized backgrounds than Eq. (5), see Ref. [9]. On the background (8), the three phenomenologically relevant sectors feel the magnetic fluxes as shown in Table 1. From the table, we see the magnetic fluxes

	Left-handed $M_C^{(i)} - M_L^{(i)}$	Higgs $M_L^{(i)} - M_R^{(i)}$	Right-handed $M_R^{(i)} - M_C^{(i)}$
T_1^2	-3	6	-3
T_2^2	+1	-1	0
T_3^2	0	-1	+1

Table 1: The number of magnetic fluxes felt by each sector on each torus.

realize the three generations of quark and lepton multiplets, and the six generations of Higgs multiplets. There also remain some extra fields, e.g., massless exotics and open string moduli as mentioned above, but we can eliminate most of them with a certain orbifold projection without any changes in the MSSM sector. For such a purpose, we consider a Z_2 orbifold defined on two tori $(T^2)_2$ and $(T^2)_3$ as

$$z_1 \rightarrow z_1, \quad z_2 \rightarrow -z_2, \quad z_3 \rightarrow -z_3,$$

with the Z_2 twists acting on the fields of the 10D SYM theory as

$$\begin{aligned} V(x_\mu, z_1, z_2, z_3) &= +PV(x_\mu, z_1, -z_2, -z_3)P, \\ \phi_1(x_\mu, z_1, z_2, z_3) &= +P\phi_1(x_\mu, z_1, -z_2, -z_3)P, \\ \phi_2(x_\mu, z_1, z_2, z_3) &= -P\phi_2(x_\mu, z_1, -z_2, -z_3)P, \\ \phi_3(x_\mu, z_1, z_2, z_3) &= -P\phi_3(x_\mu, z_1, -z_2, -z_3)P, \end{aligned}$$

where P is an (8×8) projection operator ($P^2 = \mathbf{1}$) written in the following form,

$$P = \begin{pmatrix} -\mathbf{1}_4 & 0 & 0 \\ 0 & +\mathbf{1}_2 & 0 \\ 0 & 0 & +\mathbf{1}_2 \end{pmatrix}.$$

After this Z_2 projection, the remaining zero-mode contents can be expressed as follows,

$$\begin{aligned} \phi_1^{\mathcal{I}ab} &= \left(\begin{array}{cc|c|cc} \Omega_C^{(1)} & \Xi_{CC'}^{(1)} & 0 & 0 & 0 \\ \Xi_{C'C}^{(1)} & \Omega_{C'}^{(1)} & 0 & 0 & 0 \\ \hline 0 & 0 & \Omega_L^{(1)} & H_u^K & H_d^K \\ \hline 0 & 0 & 0 & \Omega_{R'R'}^{(1)} & \Xi_{R'R''}^{(1)} \\ 0 & 0 & 0 & \Xi_{R''R'}^{(1)} & \Omega_{R''}^{(1)} \end{array} \right), \\ \phi_2^{\mathcal{I}ab} &= \left(\begin{array}{cc|c|cc} 0 & 0 & Q^I & 0 & 0 \\ 0 & 0 & L^I & 0 & 0 \\ \hline 0 & 0 & 0 & 0 & 0 \\ \hline 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \end{array} \right), \quad \phi_3^{\mathcal{I}ab} = \left(\begin{array}{cc|c|cc} 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \\ \hline 0 & 0 & 0 & 0 & 0 \\ \hline U^J & N^J & 0 & 0 & 0 \\ D^J & E^J & 0 & 0 & 0 \end{array} \right), \end{aligned}$$

where 5×5 block submatrices are matrix representations of the remaining gauge subgroups similar to Eq. (7). Indices $I, J = 1, 2, 3$ label the degenerate zero-modes of the matter sector, that is the three generations of quarks and leptons. The Higgs multiplets also have an index $K = 1, 2, \dots, 6$. Only a few of extra fields survive, Ω and Ξ . Again phenomenologically a certain prescription would be required to give them heavy masses somehow to be decoupled from the MSSM, which is beyond the scope of this paper.

Under the above Z_2 twist, nonvanishing Wilson-lines on tori $(T^2)_2$ and $(T^2)_3$, $\zeta_a^{(2)}, \zeta_a^{(3)} \neq 0$, are forbidden. However, only $\zeta_a^{(1)}$ is effective for the flavor structure because the magnetic fluxes shown in Eq. (8) cause the three-generation structure solely on torus $(T^2)_1$. Therefore only the Wilson-line parameters $\zeta_a^{(1)}$ are enough to control the hierarchical structure of the Yukawa matrices in the MSSM sector. We will assign specific values to $\zeta_a^{(1)}$ in the next section.

On this non-trivial background, we can derive the 4D effective action of the 10D SYM theory with the superfield description (3) and (4). It contains the gauge and matter kinetic terms, gauge interaction terms and Yukawa coupling terms of the MSSM, and the coefficients of them are given as functions of the 10D gauge coupling constant g , the torus radii $R^{(i)}$ and the torus complex structures $\tau^{(i)}$. The values of them will be given as VEVs of moduli fields in supergravity (SUGRA). The moduli supermultiplets on the magnetized tori consist of dilaton,

Kähler moduli and complex structure moduli chiral supermultiplets¹, which are denoted as superfields S , T_i and U_i respectively. We parameterize their VEVs as

$$\langle S \rangle = s + \theta\theta F^s, \quad \langle T_i \rangle = t_i + \theta\theta F^{t_i}, \quad \langle U_i \rangle = u_i + \theta\theta F^{u_i},$$

and the relations between $\{s, t_i, u_i\}$ and $\{g, R^{(i)}, \tau^{(i)}\}$ are given by

$$\text{Re } s = g^{-2} \mathcal{A}^{(1)} \mathcal{A}^{(2)} \mathcal{A}^{(3)}, \quad \text{Re } t_i = g^{-2} \mathcal{A}^{(i)}, \quad u_i = i \tau^{(i)}.$$

Each F^m of the supermultiplets $m = s, t_i, u_i$ is a parameter describing the magnitude of the SUSY breaking mediated by each moduli chiral multiplet. Using these relations, we can determine the moduli dependence of the 4D effective action and construct a 4D effective SUGRA action. The general action of the 4D $\mathcal{N} = 1$ effective SUGRA (on the 4D gravitational background) is given in the superspace as

$$S = \int d^4x \sqrt{-g^C} \left[-3 \int d^4\theta C C^* (e^{K/3}) + \left\{ \int d^2\theta \left(\frac{1}{4} f_a W^{a\alpha} W_\alpha^a + C^3 W + \right) + h.c. \right\} \right],$$

using a chiral compensator superfield C whose F -components will be denoted by F^C later. Finally, we can find the MSSM sector there, where the Kähler and superpotential, K and W , can be expanded as

$$\begin{aligned} K &= K_0 + Z_{I\bar{I}}^{(Q_L)} Q_L^I \bar{Q}_L^{\bar{I}} + Z_{J\bar{J}}^{(Q_R)} Q_R^J \bar{Q}_R^{\bar{J}} + Z_{K\bar{K}}^{(H)} H^K \bar{H}^{\bar{K}} + \dots, \\ W &= W_0 + \lambda_{IJK}^{(Q_R)} Q_L^I Q_R^J H^K + \dots. \end{aligned}$$

The chiral superfields $Q_L^I = \{Q^I, L^I\}$, $Q_R^J = \{U^J, D^J, N^J, E^J\}$ and $H^K = \{H_u^K, H_d^K\}$ as well as (implicit) vector superfields in the MSSM appear with the Kähler metrics $Z_{I\bar{I}}^{(Q_L)}$, $Z_{J\bar{J}}^{(Q_R)}$ and $Z_{K\bar{K}}^{(H)}$, the holomorphic Yukawa couplings $\lambda_{IJK}^{(Q_R)}$ as well as the gauge kinetic functions f_a as functions of the seven moduli chiral superfields. The Kähler metrics are given by

$$\begin{aligned} Z_{I\bar{I}}^{(Q_L)} &= \delta_{I\bar{I}} \frac{1}{\sqrt{3}} (T_2 + \bar{T}_2)^{-1} (U_1 + \bar{U}_1)^{-1/2} (U_2 + \bar{U}_2)^{-1/2} \exp \frac{4\pi (\text{Im } \zeta_{Q_L})^2}{3 (U_1 + \bar{U}_1)}, \\ Z_{J\bar{J}}^{(Q_R)} &= \delta_{J\bar{J}} \frac{1}{\sqrt{3}} (T_3 + \bar{T}_3)^{-1} (U_1 + \bar{U}_1)^{-1/2} (U_3 + \bar{U}_3)^{-1/2} \exp \frac{4\pi (\text{Im } \zeta_{Q_R})^2}{3 (U_1 + \bar{U}_1)}, \\ Z_{K\bar{K}}^{(H)} &= \delta_{K\bar{K}} \sqrt{6} (T_1 + \bar{T}_1)^{-1} \left\{ \prod_{i=1}^3 (U_i + \bar{U}_i)^{-1/2} \right\} \exp \frac{-4\pi (\text{Im } \zeta_H)^2}{6 (U_1 + \bar{U}_1)}, \end{aligned} \quad (10)$$

where the Wilson-lines $\zeta_{Q_L, Q_R, H}$ distinguish the up/down sectors of quark/lepton multiplets, but they are universal for the generations, which are expressed as

$$\begin{aligned} \zeta_Q &\equiv \zeta_C^{(1)} - \zeta_L^{(1)}, & \zeta_L &\equiv \zeta_{C'}^{(1)} - \zeta_L^{(1)}, & \zeta_{H_u} &\equiv \zeta_L^{(1)} - \zeta_{R'}^{(1)}, & \zeta_{H_d} &\equiv \zeta_L^{(1)} - \zeta_{R''}^{(1)}, \\ \zeta_U &\equiv \zeta_{R'}^{(1)} - \zeta_C^{(1)}, & \zeta_D &\equiv \zeta_{R''}^{(1)} - \zeta_C^{(1)}, & \zeta_N &\equiv \zeta_{R'}^{(1)} - \zeta_{C'}^{(1)}, & \zeta_E &\equiv \zeta_{R''}^{(1)} - \zeta_{C'}^{(1)}. \end{aligned}$$

¹ The definitions of moduli multiplets are the same as those for pure factorizable tori regardless of the existence of YM fluxes.

The holomorphic Yukawa couplings are written as

$$\lambda_{IJK}^{(Q_R)} = \sum_{m=1}^6 \delta_{I+J+3(m-1),K} \vartheta \left[\frac{3(I-J)+9(m-1)}{54} \right] (3(\bar{\zeta}_{Q_L} - \bar{\zeta}_{Q_R}), 54iU_1), \quad (11)$$

where the explicit generation dependences arise as a consequence of the localized zero-modes, and the Wilson lines distinguish the right-handed sectors Q_R . ϑ is a Jacobi-theta function. Finally, the gauge kinetic functions $f_a = S$ are universal for all the gauge groups. We have successfully obtained the 4D effective SUGRA action focusing on the MSSM contents.

3 Numerical analysis of Yukawa matrices

In this section, we estimate the masses and mixing angles of the quarks and leptons numerically. We have the six generations of Higgs multiplets and identify one linear combination of them as the MSSM Higgs multiplets. Then the physical Yukawa matrices are expressed as

$$\begin{aligned} y_{IJ}^u v_u &= \frac{\lambda_{IJK}^{(U)} \langle H_u^K \rangle}{\sqrt{e^{-K_0} Z^{(Q)} Z^{(U)} Z^{(H_u)}}}, & y_{IJ}^d v_d &= \frac{\lambda_{IJK}^{(D)} \langle H_d^K \rangle}{\sqrt{e^{-K_0} Z^{(Q)} Z^{(D)} Z^{(H_d)}}}, \\ y_{IJ}^\nu v_u &= \frac{\lambda_{IJK}^{(N)} \langle H_u^K \rangle}{\sqrt{e^{-K_0} Z^{(L)} Z^{(N)} Z^{(H_u)}}}, & y_{IJ}^e v_d &= \frac{\lambda_{IJK}^{(E)} \langle H_d^K \rangle}{\sqrt{e^{-K_0} Z^{(L)} Z^{(E)} Z^{(H_d)}}}, \end{aligned}$$

where v_u and v_d are the VEVs of up- and down-type Higgs fields in the MSSM respectively, and we note that there is a summation over the generations $K = 1, 2, \dots, 6$ of Higgs fields in each matrix. The denominators appear as a consequence of the canonical normalizations of fields. The moduli Kähler potential K_0 of the three factorizable tori $\prod_{i=1}^3 (T^2)_i$ is given by

$$K_0 = -\ln(S + \bar{S}) - \sum_i \ln(T_i + \bar{T}_i) - \sum_i \ln(U_i + \bar{U}_i).$$

Now we can estimate the values of the Yukawa coupling constants depending on the moduli VEVs, the Higgs VEVs and the Wilson lines.

In Ref. [6], a semi-realistic pattern of the masses and mixing angles of the quarks and leptons was realized with certain values of them shown in Appendix A. The top quark mass obtained there is slightly below the experimental data, that is, the value of the top Yukawa coupling y_{33}^u is a bit small, because such a high-precision analysis was not required for the purpose in Ref. [6] to study the flavor structure. However, large quantum corrections are required to realize the 126 GeV Higgs mass within the MSSM (or MSSM-like models) and the dominant contribution will come from the top Yukawa coupling y_{33}^u . The slightly small top Yukawa coupling induced by the ansatz adopted in Ref. [6] will be a disadvantage to obtain a realistic model including the 126 GeV Higgs boson.

Therefore, for the purpose in this paper to study especially the Higgs mass, we adopt another ansatz a little different from the original one. Our new ansatz is the following; the VEVs of

Higgs fields are

$$\begin{aligned}
\tan \beta &\equiv v_u/v_d = 15, \\
\langle H_u^K \rangle &= (0.0, 0.0, 3.3, 1.2, 0.0, 0.0) \times v_u \mathcal{N}_u, \\
\langle H_d^K \rangle &= (0.0, 0.1, 5.9, 5.9, 0.0, 0.1) \times v_d \mathcal{N}_d,
\end{aligned} \tag{12}$$

where we use the normalization factors $\mathcal{N}_u = \sqrt{3.3^2 + 1.2^2}$ and $\mathcal{N}_d = \sqrt{2(0.1^2 + 5.9^2)}$. The VEVs of moduli are chosen as

$$\begin{aligned}
\pi s &= 6.0, \\
(t_1, t_2, t_3) &= (3.0, 1.0, 1.0) \times 2.8 \times 10^{-8}, \\
(u_1, u_2, u_3) &= (4.4, 1.0, 1.0).
\end{aligned} \tag{13}$$

The dilaton VEV s leads to $4\pi/g_a^2 = 24$ at the GUT scale $M_{\text{GUT}} = 2.0 \times 10^{16}$, which is the unified value implemented by the MSSM². The VEVs t_r ($r = 1, 2, 3$) of Kähler moduli determine the compactification scale, which is fixed to M_{GUT} , and the ratios between them are chosen to satisfy the SUSY condition (9). As for the VEVs u_r ($r = 1, 2, 3$) of complex structure moduli, only the first u_1 is important because the flavor structure originates solely from the first torus $(T^2)_1$. The other u_2 and u_3 have no effect on the flavor structures and we set them as $u_2 = u_3 = 1.0$ for simplicity.

The Wilson-lines, which mostly affect the Yukawa hierarchy, are selected as

$$\left(\zeta_C^{(1)}, \zeta_{C'}^{(1)}, \zeta_L^{(1)}, \zeta_{R'}^{(1)}, \zeta_{R''}^{(1)} \right) = (0.0, -0.5i, -0.6i, 0.9i, 0.8i). \tag{14}$$

We have a degree of freedom to shift the wavefunctions of all the elements universally, that allows us to set $\zeta_C^{(1)} = 0$ for simplicity without affecting the flavor structure. The above ansatz yields a semi-realistic pattern of the quark masses, the charged lepton masses and the CKM mixing angles [10] at the electroweak (EW) scale M_Z including contributions from the full 1-loop Renormalization Group Equation (RGE) which are shown in Table 2. We obtain the larger top quark mass in spite of the smaller value of $\tan \beta$. The top Yukawa coupling $y_{33}^u = 0.997$ becomes larger than $y_{33}^u = 0.971$ obtained in Ref. [6] at the EW scale. That causes an unnegligible effect because the dominant correction to the Higgs mass is proportional to $(y_{33}^u)^4$. Moreover, the smaller value of $\tan \beta$ has another advantage in the next section.

Note that, as in Ref. [6], we assume the existence of supersymmetric Higgs mass term which could not be derived from the 10D SYM action. Non-perturbative effects and/or higher-dimensional operators might be able to generate the mass term,

$$\mu_{KL} H_u^K H_d^L,$$

in the superpotential W of the 4D effective theory. This μ_{KL} is a (6×6) matrix and, as mentioned previously, we assume that the five eigenvalues are as large as M_{GUT} and the last

² Note again that all the exotics other than MSSM contents are assumed to be heavier than the compactification scale

	Sample values	Observed
(m_u, m_c, m_t)	$(2.43 \times 10^{-3}, 0.431, 1.73 \times 10^2)$	$(2.30 \times 10^{-3}, 1.28, 1.74 \times 10^2)$
(m_d, m_s, m_b)	$(4.59 \times 10^{-3}, 1.86 \times 10^{-1}, 10.7)$	$(4.8 \times 10^{-3}, 0.95 \times 10^{-1}, 4.18)$
(m_e, m_μ, m_τ)	$(1.53 \times 10^{-3}, 6.36 \times 10^{-2}, 5.11)$	$(5.11 \times 10^{-4}, 1.06 \times 10^{-1}, 1.78)$
$ V_{\text{CKM}} $	$\begin{pmatrix} 0.987 & 0.161 & 0.00585 \\ 0.159 & 0.982 & 0.0964 \\ 0.0213 & 0.0942 & 0.995 \end{pmatrix}$	$\begin{pmatrix} 0.97 & 0.23 & 0.0035 \\ 0.23 & 0.97 & 0.041 \\ 0.0087 & 0.040 & 1.0 \end{pmatrix}$

Table 2: The sample theoretical values of the quark masses, the charged lepton masses and the CKM mixing angles, which are estimated at the EW scale through the full 1-loop RGE flows. The observed values are quoted from Ref. [11].

one, which will be identified with the so-called μ -parameter of the MSSM, is much smaller than the other five comparable to the (low) SUSY breaking scale. The five linear combinations among the six generations of the pair of Higgs doublets (H_u^K, H_d^K) are decoupled from the MSSM at above the GUT scale and the last one pair is identified as the pair of MSSM Higgs doublets (H_u, H_d) having the above mentioned μ -term, $\mu H_u H_d$, in the superpotential³.

As for the neutrino sector, we can study it assuming the seesaw mechanism [12] with the following heavy Majorana mass term,

$$M^N = \begin{pmatrix} 0.1 & 1.9 & 0.0 \\ 1.9 & 0.3 & 3.1 \\ 0.0 & 3.1 & 1.4 \end{pmatrix} \times 10^{11} \text{ GeV},$$

which might be also generated by non-perturbative and/or higher-order effects. In this case, a semi-realistic pattern of the neutrino masses and the PMNS mixing angles [13], as shown in Table 3, are produced at the same time as the quark and the charged lepton masses as well as the CKM mixing angles shown in Table 2 are realized.

In this section, we have shown a sample spectrum of the SM particles with the ansatz (12)-(14). On this background, we will study the SUSY spectra and the Higgs boson mass in the next section. Finally, we emphasize that the observed mysterious hierarchical structure of quarks and leptons are successfully generated from the non-hierarchical input values of VEVs shown in Eqs. (12)-(14), by virtue of the wavefunction localization of chiral zero-modes caused by magnetic fluxes.

³ We can also study with another scenario in which the five pairs of Higgs doublets other than the MSSM Higgs doublets are not decoupling from the MSSM at the GUT scale and would affect on the phenomenologies. We will study that elsewhere.

	Sample values	Observed
$(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$	$(2.57 \times 10^{-17}, 1.11 \times 10^{-11}, 9.30 \times 10^{-11})$	$< 2 \times 10^{-9}$
$ m_{\nu_1}^2 - m_{\nu_2}^2 $	1.24×10^{-22}	7.50×10^{-23}
$ m_{\nu_1}^2 - m_{\nu_3}^2 $	8.65×10^{-21}	2.32×10^{-21}
$ V_{\text{PMNS}} $	$\begin{pmatrix} 0.933 & 0.255 & 0.255 \\ 0.354 & 0.779 & 0.518 \\ 0.0668 & 0.574 & 0.816 \end{pmatrix}$	$\begin{pmatrix} 0.82 & 0.55 & 0.16 \\ 0.51 & 0.58 & 0.64 \\ 0.26 & 0.61 & 0.75 \end{pmatrix}$

Table 3: The sample theoretical values of the neutrino masses and the PMNS mixings, which are estimated at the EW scale through the full 1-loop RGE flows. The observed values are quoted from Ref. [11].

4 The Higgs boson mass and SUSY spectra

We study the Higgs boson mass and SUSY spectra in the 4D effective SUGRA. We consider two types of SUSY breaking mediation mechanisms in this model, the moduli (S , T_r , U_r) and the anomaly (compensator C) [14] mediations whose mixture is referred to as mirage mediation [15, 16]. The soft SUSY breaking parameters induced by them are calculated by the formulae given by Ref. [15] and we estimate the spectra at the EW scale through the full 1-loop RGE flows with the MSSM contents numerically, and then we can study the Higgs boson mass.

4.1 Soft parameters

In this section, we assume a certain moduli stabilization mechanism works and it determines the VEVs of seven moduli scalar components, s, t_r, u_r , which is consistent with Eqs. (12)-(14) to realize the SM flavors. Our model has two types of SUSY breaking mediations and, although their contributions proportional to the mediator's F -terms will also be determined by the moduli stabilization as well as the SUSY breaking scenario, we study here treating the VEVs F^m ($m = s, t_r, u_r$) and F^C as free parameters representing the magnitude of the SUSY breaking mediated by each of the moduli and the anomaly respectively. The nonvanishing F -terms generate the soft SUSY breaking terms, and then, the soft parameters can be calculated and the spectrum is obtained depending on these parameters. In this way, we will be able to study the model concentrating on the Higgs boson mass and the SUSY spectrum without a concrete moduli stabilization scenario. Inversely speaking, the results of our analysis would probe the moduli stabilization as well as the SUSY breaking mechanisms behind our model.

As discussed above, we can calculate the soft parameters, which are the gaugino masses M_a , the scalar tri-linear couplings called A -terms a_{IJK} and the soft scalar mass squares $(m_Q^2)_I^J$, by

	S	T_1	T_2	T_3	U_1	U_2	U_3
H	1/3	-2/3	1/3	1/3	-1/6	-1/6	-1/6
Q_L	1/3	1/3	-2/3	1/3	-1/6	-1/6	1/3
Q_R	1/3	1/3	1/3	-2/3	-1/6	1/3	-1/6

Table 4: The values of c_Q^m appear in Eq.(15).

using the following formulae [15],

$$\begin{aligned}
M_a &= \frac{F^s}{s + \bar{s}} + \frac{b_a}{16\pi^2} g_a^2 \frac{F^C}{C}, \\
a_{IJK} &= y_{IJK}(c_{Q_L}^m + c_{Q_R}^m + c_H^m) \frac{F^m}{\varphi_m + \bar{\varphi}_m} + y_{IJK} F^m \partial_m \ln \lambda_{IJK} + (y_{LJK} \gamma_I^L + (I \leftrightarrow J, K)) \frac{F^C}{C}, \\
(m_{\tilde{Q}}^2)_I^J &= c_Q^m \left| \frac{F^m}{\varphi_m + \bar{\varphi}_m} \right|^2 \delta_I^J - \partial_m \gamma_I^J \left(\frac{F^m}{\varphi_m + \bar{\varphi}_m} \frac{\bar{F}^C}{C} + \text{h.c.} \right) + \frac{1}{4} \dot{\gamma}_I^J \left| \frac{F^C}{C} \right|^2, \tag{15}
\end{aligned}$$

where γ_I^J is the anomalous dimension and $\dot{\gamma} = \frac{\partial \gamma}{\partial \ln \mu/\Lambda}$. Note that, the index "m" is summed over the seven moduli supermultiplets (S , T_r , U_r), and $\varphi_m = s, t_r, u_r$ represents the VEVs of their scalar components. The values of c_Q^m ($Q = Q_L, Q_R, H$) derived from the Kähler metrics (10) are listed in Table 4. They are universal for their generations and determine how each moduli contributes to the soft masses of the left-handed Q_L , right-handed Q_R and Higgs H scalars.

From Table 4, we see that some of them give negative contributions. Negative contributions to the soft mass squares are disfavored in at least the left-handed and right-handed sectors to avoid tachyons and we expect that the total contributions given by the seven moduli should be positive⁴. For instance, the five moduli other than S and T_1 give the negative contributions to either the left or the right, or both. Thus one of them cannot solely contribute to the soft masses. We also find that the net contributions made by the three Kähler moduli T_r vanish if the three contributions are equal, $F^{t_1} = F^{t_2} = F^{t_3}$. The gaugino masses are generated by dilaton S . We expect a naively nonvanishing F^s is required to some extent to generate gluino masses large enough satisfying the experimental lower bound. The moduli contributions to the A -terms are determined by the sums $c_{Q_L}^m + c_{Q_R}^m + c_H^m$, then we find $c_{Q_L}^m + c_{Q_R}^m + c_H^m \neq 0$ for $m = S, U_1$ from Table 4 and the other moduli will not affect the A -terms. Moreover, we also find $c_{Q_L}^S + c_{Q_R}^S + c_H^S = 1$ and that is equivalent to the so-called mirage condition [15] if F^{u_1} is negligible, which is required to avoid the SUSY flavor violations as shown later.

The moduli dependence of the soft terms (15) is determined by the configuration of magnetic fluxes realizing the SM flavor structures. The Kähler modulus T_i appears in the Kähler metrics of the chiral multiplets embedded in ϕ_i defined in Eq. (2) and it contains bosons behaving as vectors on only the i -th torus $(T^2)_i$ and their partners. For example, $Z_{II}^{(Q_L)}$ depends on the

⁴ Even if the anomaly mediation is included, we need the positive contribution somewhat because the pure anomaly mediation generally induces tachyonic sleptons.

modulus T_2 as shown in Eq. (10) because the left-handed matters Q^I and L^I are embedded into ϕ_2 . On the other hand, as for the complex structure moduli dependence, the modulus U_i appears in the Kähler metrics if the corresponding multiplets feel nonvanishing magnetic fluxes on the i -th torus $(T^2)_i$. We see from Table 1 that all the three sectors have nonvanishing fluxes on the first torus $(T^2)_1$, consequently, the complex structure modulus U_1 appears in their Kähler metrics. And also, the left-handed sector feels no magnetic fluxes on the third torus $(T^2)_3$, for example, and the modulus U_3 will not appear in its Kähler metric.

The modulus U_1 will receive the other severe constraint. The moduli dependence of all the Kähler metrics is universal for each generation involved in, while the modulus U_1 appears in the holomorphic Yukawa couplings (11) depending on the generation indices, because the three-generation structures are caused by the magnetic fluxes on the first torus $(T^2)_1$. That induces a severe restriction on the magnitude of the SUSY braking mediated by the complex structure modulus U_1 to suppress the SUSY flavor violations.

A phenomenological analysis of this model has in part been done in Ref. [6] neglecting Yukawa matrix elements other than the most dominant one $y_{33}^{u,d,\nu,e}$. In this paper, we include the contributions from all of the Yukawa couplings to soft parameters at the GUT scale and their 1-loop RG effects, and we evaluate superparticle masses and SUSY flavor violations more precisely taking the latest experimental data into account. Furthermore, we estimate the Higgs boson mass, which has never estimated in this model, by calculating the 1-loop effective potential containing the top and bottom (s)quarks corrections [17]⁵.

There are only the eight remaining parameters undetermined after realizing the semi-realistic SM sector. They are the VEVs F^m ($m = s, t_r, u_r$) and F^C of the F -terms of the moduli and compensator chiral supermultiplets, respectively. Since only the dilaton S appears in the gauge kinetic functions and can give the masses to gauginos at the tree level, we refer to the normalized F^s as the overall SUSY breaking scale,

$$M_{\text{SB}} = \sqrt{K_{\bar{S}S}} F^s,$$

and parameterize the other contributions by the following normalized ratios to M_{SB} ,

$$R_r^T = \frac{\sqrt{K_{\bar{T}_r T_r}} F^{t_r}}{M_{\text{SB}}}, \quad R_r^U = \frac{\sqrt{K_{\bar{U}_r U_r}} F^{t_r}}{M_{\text{SB}}}, \quad R^C = \frac{1}{\ln M_p/m_{3/2}} \frac{F^C/C}{M_{\text{SB}}}.$$

We show the results of numerical analyses varying these parameters in the following.

4.2 Numerical results

First, we study M_{SB} dependence of the Higgs boson mass without the other moduli contributions, $R_1^T = R_2^T = R_3^T = R_1^U = R_2^U = R_3^U = 0$. This situation corresponds to the simplest single modulus scenario and there remain two parameters, M_{SB} and R^C . We exhibit the contours of the Higgs boson mass and some observationally relevant curves (regions) on the $(M_{\text{SB}} [\text{TeV}], R^C)$ -plane and show that in Fig. 1. In this figure, the theoretical value of the Higgs boson mass is

⁵ In Ref. [17], the leading log approximation was used, but we evaluate the self-coupling constant of the Higgs boson solving RGEs numerically for a better accuracy.

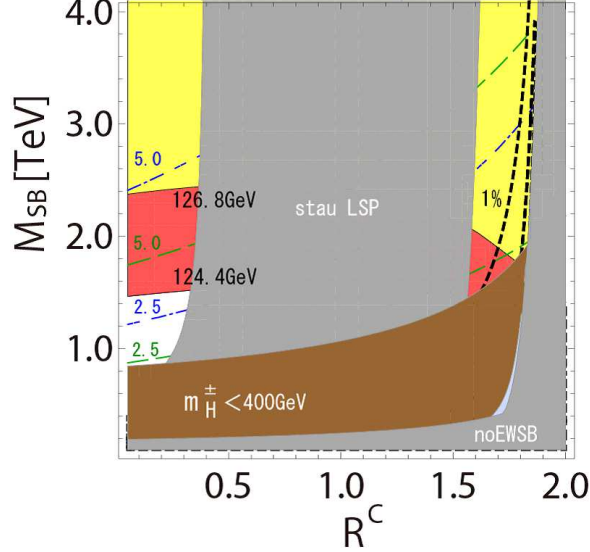


Figure 1: The contours of the Higgs boson masses and some experimentally relevant curves (regions) are drawn on the $(M_{\text{SB}} [\text{TeV}], R^C)$ -plane with $R_1^T = R_2^T = R_3^T = R_1^U = R_2^U = R_3^U = 0$. In the red regions, the theoretical value of the Higgs boson mass resides in the allowed range, $124.4 < m_h < 126.8$ GeV. The Higgs boson is heavier than that in the yellow regions.

inside the range of the experimental observations [1, 2], $124.4 < m_h < 126.8$ GeV, in the red regions and exceeds that in the yellow regions. We find the 126 GeV Higgs boson can be realized with $M_{\text{SB}} \sim 2.0$ TeV. In the gray regions, the EW symmetry breaking (EWSB) will not occur successfully or a stau becomes the Lightest Supersymmetric Particle (LSP), with which the pure MSSM cannot give any candidate for the dark matter. The green and blue dashed lines represent the masses (TeV) of the gluino and the lighter top squark respectively. The black dashed lines correspond to the degree of tuning the higgsino mass parameter (μ -parameter) to obtain the observed Z -boson mass (radiatively), which is defined as $100/\Delta_\mu$ (%) with

$$\Delta_\mu = \left| \frac{\partial \log m_Z^2}{\partial \log \mu^2} \right|.$$

The whole parameter space of (M_{SB}, R^C) shown in Fig. 1 is free from the experimental constraints on various SUSY flavor violations estimated by evaluating the mass insertion parameters [18]. This is mostly because here we set $R_1^U = 0$, and the effect of $R_1^U \neq 0$ will be shown later. We should remark that the lower bound on the charged Higgs boson mass is treated as $m_{H^\pm} > 400$ GeV, because processes with charged Higgs boson exchange would contribute to $\Gamma(b \rightarrow s\gamma)$ and the charged Higgs boson lighter than 350 GeV is disfavored⁶ [19].

⁶ The analysis with the mass insertion parameters cannot take into account this contribution, because the corresponding Feynman diagrams do not contain the soft parameters.

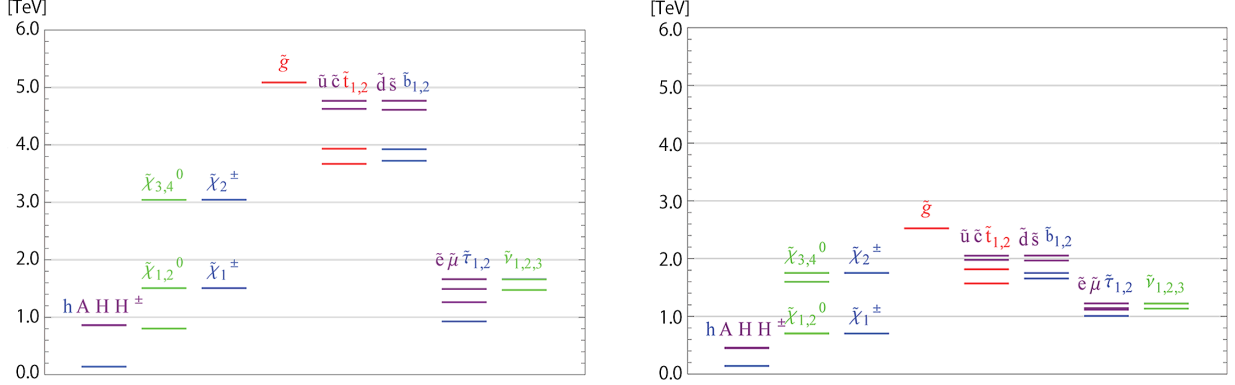


Figure 2: The mass spectra of the MSSM particles with $M_{\text{SB}} = 1.8$ TeV and $R^C = 0.1$ (1.7) in the left (right) panel.

We exclude the region, where $m_{H^\pm} < 400$ GeV, using brown shade.

Thus the suitable value of M_{SB} is about 2.0 TeV for the 126 GeV Higgs boson in this case, with which the SUSY spectra consistent with the various experimental results can be also obtained. Furthermore, the fine-tuning is relaxed better than 1% with $R^C \sim 1.7$ and then almost all the soft parameters are unified at around TeV scale as pointed out in the TeV scale mirage mediation models [20] (See also Refs. [21, 22]). Since the mirage unification at the TeV scale leads to the relatively light charged Higgs bosons as long as $|\mu|^2$ is small, consequently, some of the natural region $100/\Delta_\mu \gtrsim 1$ % is covered by the brown shade, where the charged Higgs boson is lighter than 400 GeV because of the mirage unification, but some natural regions still remain allowed.

We remark that, with the small top Yukawa ansatz of Ref. [6], the 126 GeV Higgs mass requires the higher SUSY breaking scale than 2.0 TeV, indeed, we could not find the allowed region with $M_{\text{SB}} = 2.0$ TeV. Furthermore, the small top Yukawa coupling will be accompanied with the large value of $\tan \beta$. In general, a large $\tan \beta$ induces the light charged Higgs boson and then broadens the excluded brown shade region in Fig. 1, which will cover the natural region around $R^C \sim 1.7$. Thus the new Yukawa ansatz is more favored in order to realize the Higgs boson mass without the fine-tuning.

We calculate the sample theoretical SUSY spectra given at the two different points in Fig. 1. We show the spectrum derived from $(M_{\text{SB}}, R^C) = (1.8 \text{ TeV}, 0.1)$ in the left panel of Fig. 2 and the other one from $(1.8 \text{ TeV}, 1.7)$ in the right. In the former (left) case, a CMSSM-like spectrum is realized with the 0.048 % tuning and the 125.4 GeV Higgs boson. In this case, the LSP is a bino-like neutralino and the colored particles are heavier than the non-colored particles by sub-TeV. We can expect that sleptons or electroweakinos will be discovered earlier. In the latter (right) case, we get the 126.5 GeV Higgs boson with the relaxed tuning, 1.1 %. The LSP is a higgsino-like neutralino, and most of the other sparticles are lighter than about 2.0 TeV and can be reached at the LHC in the near future. These two spectra carry the different LSP. We can further study this model with results of the cosmological observations if we consider the LSP dark matter, but we will not execute it here just leaving some comments. The bino dark

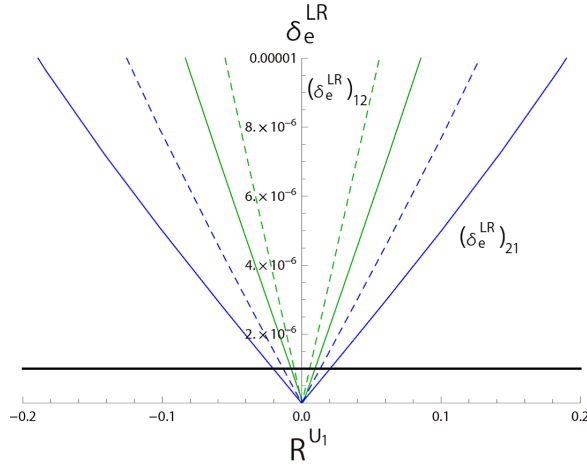


Figure 3: The mass insertion parameters $(\delta_e^{LR})_{12}$ (green) and $(\delta_e^{LR})_{21}$ (blue) as functions of R_1^U with $R^C = 0.1$ (solid) or $R^C = 1.7$ (dashed). The black horizontal line represents $\delta_e^{LR} = 10^{-6}$.

matter scenario is being desperate as known. In the other case with a higgsino-like neutralino LSP, we can expect there are light colored SUSY particles and their masses are bounded by the Higgs observation because the theoretical value of the Higgs boson mass exceeds the observed one in the yellow regions in Fig. 1 (where the value of $\tan \beta$ is fixed). We might be able to verify this model by combination of the various experiments (at least with the parameters shown in Section 3).

As mentioned above, the SUSY flavor violations are mostly depending on the complex structure modulus U_1 of the first torus $(T^2)_1$, and here we will study its effects. Within $M_{\text{SB}} \leq 1.0$ TeV, the modulus U_1 must not participate in the SUSY breaking mediation to suppress the dangerous SUSY flavor violations [6], especially, concerning the process $\mu \rightarrow e\gamma$. We study the effect of $R_1^U \neq 0$ more precisely with the above two sample spectra. We show the relevant mass insertion parameters as functions of R_1^U in Fig. 3 where $M_{\text{SB}} = 1.8$ TeV and $R^C = 0.1$ (1.7) with the solid (dashed) lines, and the others are vanishing. They have the stringent constraints $\mathcal{O}((\delta_e^{LR})_{12,21}) \lesssim 10^{-6}$ (black horizontal line) given by Ref. [23]. We see that the value of R_1^U must be tiny, at least $|R_1^U| \lesssim 0.01$, even if the SUSY spectra are a little heavier to be consistent with the Higgs discovery. We can also satisfy all the other constraints on SUSY flavor violations from the flavor changing neutral current (FCNC) experiments easily with the tiny value of $|R_1^U|$. We adopt $R_1^U = 0$ for simplicity in the following analysis⁷.

Next, we study our model with the other moduli-mediated contributions. The moduli T_2 , T_3 , U_2 and U_3 of the other tori $(T^2)_2$ and $(T^2)_3$ than $(T^2)_1$ can solely give negative contributions to the squark and slepton mass squares as is seen from Table 4. As mentioned before, the totally positive contributions are required in the left- and right-handed sectors to obtain non-tachyonic sparticles. (Furthermore, if the absolute values of the squared soft masses become so large in the Higgs sector, the fine-tuning problem can be serious.) We consider the

⁷ The KKLT-type scenario [24] of moduli stabilization, fixing all the complex structure moduli at a high-scale in a supersymmetric way, is one of the candidates for realizing $R_1^U \approx 0$.

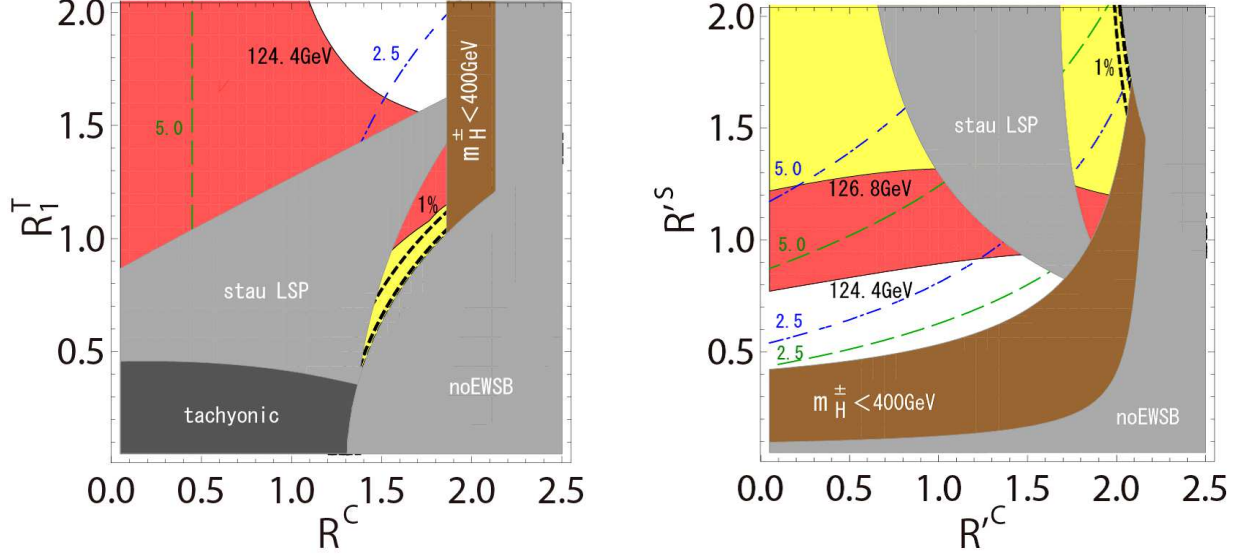


Figure 4: The contours of the Higgs boson mass on the (R_1^T, R^C) -plane with $M_{\text{SB}} = 2.0$ TeV and $R_2^T = R_3^T = 1$ in the left panel. We also represent various observational constraints similar to those in Fig. 1. The similar drawing is shown on the (R'^S, R'^C) -plane with $M'_{\text{SB}} = 2.0$ TeV in the right panel.

case with $R_2^T = R_3^T$ and $R_2^U = R_3^U$ for simplicity. A combination of U_2 and U_3 behaves similarly to T_1 as long as we focus on the SUSY spectra, and we can choose $R_2^U = R_3^U = 0$ without loss of generality.

We show again the contours of the Higgs boson mass and some observationally relevant curves (regions) on the (R_1^T, R^C) -plane with $M_{\text{SB}} = 2.0$ TeV and $R_2^T = R_3^T = 1$ in the left panel of Fig. 4. At the points with $R_1^T \sim 1$, the net contributions from the three Kähler moduli T_r almost vanish and dilaton S dominates the SUSY breaking mediation. The negative contributions from T_2 and T_3 dominate below the points, $R_1^T < 1$, while the positive one from T_1 becomes dominant above them, $R_1^T > 1$. We can see the negative contributions comparable to that of the dilaton are dangerous, and

$$\sum_{m \neq S} c_Q^m R^m \gtrsim 0, \quad \text{for} \quad Q = Q_L, Q_R,$$

is required in general cases.

Since some observationally excluded regions will shrink with $R_1^T \gtrsim 1$, we further study the T_1 -dominant case in another analysis, where we renew some parameters as follows,

$$M'_{\text{SB}} = \sqrt{K_{\bar{T}_1 T_1}} F^{t_1} = 2.0 \text{ TeV}, \quad R'^S = \frac{\sqrt{K_{\bar{S} S}} F^s}{M'_{\text{SB}}}, \quad R'^C = \frac{1}{\ln M_p / m_{3/2}} \frac{F^C / C}{M'_{\text{SB}}},$$

and the other F -terms are vanishing. Similarly we show the contours of the Higgs boson mass and some experimentally relevant curves (regions) on the (R'^S, R'^C) -plane in the right panel of Fig. 4. Compared with Fig. 1, it is obvious that the positive contributions from T^1 to the sfermions are favored. However, we can also see the dilaton contribution is indispensable to a certain extent even if with the others contribute. The main reason of that is the large gluino mass is required, especially in the RGE flows, to yield the successful EWSB and that can be generated by only the dilaton and anomaly, because, as explained at the end of Section 2, the gauge kinetic functions are functions of only the dilaton in 10D SYM theory.

5 Conclusions and Discussions

We have studied the Higgs boson mass and precise SUSY spectra of the particle physics model derived from the 10D SYM theory compactified on the magnetized tori. This model was proposed in Ref. [6], where the magnetic fluxes in the extra dimensional space originate the complicated flavor structures of the SM and the magnitude of the SUSY flavor violations was estimated with rough approximations to check consistency. In this paper, we are focusing on the 126 GeV Higgs boson mass. For such a purpose, we have first done a minor improvement of the Yukawa structures to enhance the top Yukawa coupling y_{33}^u realizing the masses and mixing angles of the quarks and leptons. The enhanced top Yukawa coupling y_{33}^u can be an advantage to generate the large quantum corrections required to realize the 126 GeV Higgs boson within low scale SUSY breaking scenarios.

On the improved background, we have studied the model using the SUSY breaking parameters F^m and F^C of the moduli and compensator chiral multiplets, respectively. We have estimated the Higgs boson mass in the simplest case by a varying the overall SUSY breaking scale M_{SB} , and obtained the 126 GeV Higgs boson mass when the overall scale is about 2.0 TeV. In this case, the fine-tuning can be relaxed if there is a comparable contribution from the anomaly mediation with which the TeV scale mirage scenario [20] is realized. We have shown the two sample SUSY spectra allowed by the experimental constraints. One is a moduli dominated scenario and the other corresponds to the TeV scale mirage scenario. Both spectra will be reached at future experiments. They have some significant differences in the mass scales of the colored particles and the LSP constituent. These differences provide a motivation to study the model from the perspective of not only the high-energy experiments but also the cosmological observations elsewhere in a separate work. In particular, results from the combined studies might be able to verify this model in the near future.

We notice that the unimproved original flux configuration given in Ref. [6] requires the higher SUSY breaking scale, and the wide regions of the parameter space including the natural regions $100/\Delta_\mu \gtrsim 1\%$ will be excluded because of the light charged Higgs boson possibly induced by the large value of $\tan\beta$. Our new ansatz of the Yukawa matrices is more favored clearly from these points of view.

With the two sample SUSY spectra, we have estimated the magnitude of SUSY flavor violations without the approximations adopted in Ref. [6] and compared with the latest experimental data. The most stringent bound is still coming from the process $\mu \rightarrow e\gamma$ and it requires that the

complex structure modulus U_1 of the first torus $(T^2)_1$, where the three-generation structures are solely caused, should not mediate the SUSY breaking contributions, $|R_1^U| \lesssim 0.01$. In other words, U_1 modulus should be quite heavy enough to decouple from the MSSM sector in the low energy effective field theory. In such a case, all the other SUSY FCNC constraints are easily satisfied.

We have also studied the other moduli dependence of the MSSM spectra. There are some phenomenological constraints coming from other than the observed 126 GeV Higgs mass, e.g, the FCNCs, the successful (radiative) EWSB (the observed Z -boson mass), no tachyons (color and charge breaking minima) and the (LSP) dark matter candidate, etc. They restrict the dynamics of the moduli, and we find the following indications. First, dilaton S has to contribute to the SUSY breaking mediations to some extent even if the other contributions can be expected or not, because the successful EWSB requires the large gluino mass and the dilaton contribution is indispensable for that. Second, negative contributions to the squark and slepton mass squares comparable to that from the dilaton should be forbidden to avoid tachyonic particles.

We see that the various experimental results restrict the VEVs s, t_r, u_r and moduli F -terms F^s, F^{t_r}, F^{u_r} , those has their own geometrical meanings, as well as the F -term of chiral compensator F^C . The further high-energy experiments will be able to prove the SUSY breaking mediation mechanisms in our model and, furthermore, the moduli stabilization and SUSY breaking mechanisms behind it. Based on the results obtained in this paper, we will study a concrete moduli stabilization scenario including a SUSY breaking sector elsewhere. Since the recent cosmological observations is quite promoted, it would also be attractive to study about cosmological issues based on our model as mentioned in the previous section.

We can also consider some other extensions of this model including the gauge mediations. Our SUSY spectra given in this paper can be deflected in such a extended models. Although the additional gauge mediation is not expected to be a promising advantage from the naturalness perspective [22], we might be able to expect drastic changes of the spectra and be inspired to go on to more various phenomenologies.

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A Yukawa matrices of the previous work

We show, for a reference, the original input parameters adopted in Ref. [6] below,

$$\begin{aligned}
\tan \beta &= 25, \\
\langle H_u^K \rangle &= (0.0, 0.0, 2.7, 1.3, 0.0, 0.0) \times v_u \mathcal{N}_u, \\
\langle H_d^K \rangle &= (0.0, 0.1, 5.8, 5.8, 0.0, 0.1) \times v_d \mathcal{N}_d, \\
\pi s &= 6.0, \\
(t_1, t_2, t_3) &= (3.0, 1.0, 1.0) \times 2.8 \times 10^{-8}, \\
(u_1, u_2, u_3) &= (4.1, 1.0, 1.0), \\
(\zeta_C^{(1)}, \zeta_{C'}^{(1)}, \zeta_L^{(1)}, \zeta_{R'}^{(1)}, \zeta_{R''}^{(1)}) &= (0.0, 0.3i, -1.0i, 1.9i, 1.4i),
\end{aligned}$$

where the Majorana neutrino masses were assumed as

$$M^N = \begin{pmatrix} 1.1 & 1.3 & 0 \\ 1.3 & 0 & 3.2 \\ 0 & 3.2 & 1.8 \end{pmatrix} \times 10^{12} \text{ GeV}.$$

From these input parameters, the theoretical values of quark and lepton masses and mixing angles shown in Table 5 and 6 are obtained, which are compared with those shown in Table 2 and 3 derived from the improved parameters proposed in this paper.

	Sample values	Observed
(m_u, m_c, m_t)	$(3.1 \times 10^{-3}, 1.01, 1.70 \times 10^2)$	$(2.30 \times 10^{-3}, 1.28, 1.74 \times 10^2)$
(m_d, m_s, m_b)	$(2.8 \times 10^{-3}, 1.48 \times 10^{-1}, 6.46)$	$(4.8 \times 10^{-3}, 0.95 \times 10^{-1}, 4.18)$
(m_e, m_μ, m_τ)	$(4.68 \times 10^{-4}, 5.76 \times 10^{-2}, 3.31)$	$(5.11 \times 10^{-4}, 1.06 \times 10^{-1}, 1.78)$
$ V_{\text{CKM}} $	$\begin{pmatrix} 0.98 & 0.21 & 0.0023 \\ 0.21 & 0.98 & 0.041 \\ 0.011 & 0.040 & 1.0 \end{pmatrix}$	$\begin{pmatrix} 0.97 & 0.23 & 0.0035 \\ 0.23 & 0.97 & 0.041 \\ 0.0087 & 0.040 & 1.0 \end{pmatrix}$

Table 5: The sample theoretical values of the quark and charged lepton masses and the CKM mixing angles at the EW scale through the 1-loop RGE flows, those are derived from the input parameters given in Ref. [6].

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	Sample values	Observed
$(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$	$(3.6 \times 10^{-19}, 8.8 \times 10^{-12}, 2.7 \times 10^{-11})$	$< 2 \times 10^{-9}$
$ m_{\nu_1}^2 - m_{\nu_2}^2 $	7.67×10^{-23}	7.50×10^{-23}
$ m_{\nu_1}^2 - m_{\nu_3}^2 $	7.12×10^{-22}	2.32×10^{-21}
$ V_{\text{PMNS}} $	$\begin{pmatrix} 0.85 & 0.46 & 0.25 \\ 0.50 & 0.59 & 0.63 \\ 0.15 & 0.66 & 0.73 \end{pmatrix}$	$\begin{pmatrix} 0.82 & 0.55 & 0.16 \\ 0.51 & 0.58 & 0.64 \\ 0.26 & 0.61 & 0.75 \end{pmatrix}$

Table 6: The sample theoretical values of the neutrino masses and the MNS mixings at the EW scale through the 1-loop RGE flows, those are derived from the input parameters given in Ref. [6].

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